

# Spin transition in the fractional quantum Hall regime: Effect of wave function extent

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Using a magnetocapacitance technique, we determine the magnetic field of the spin transition,  $B^*$ , at filling factor  $\nu = 2/3$  in the 2D electron system in GaAs/AlGaAs heterojunctions. The field  $B^*$  is found to decrease appreciably as the wave function extent controlled by back gate voltage is increased. Our calculations show that the contributions to the shift of  $B^*$  from the  $g$  factor change due to nonparabolicity and the change of the Coulomb energy are approximately the same. The observed relative shift of  $B^*$  is described with no fitting parameters.

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Spin degrees of freedom are important in determining the ground states and excitations of the fractional quantum Hall effect in not too strong magnetic fields. In view of the competition between the Coulomb energy in the two-dimensional (2D) electron system,  $E_C \propto e^2/\kappa l_B$  (where  $\kappa$  is the dielectric constant and  $l_B = (\hbar c/eB)^{1/2}$  is the magnetic length), and the Zeeman energy,  $E_Z = g\mu_B B$ , the fully polarized states in the high-field limit should become depolarized with decreasing magnetic field. This manifests itself as spin transitions between fully and partially polarized states and between states with different partial polarization. The ratio of the Zeeman and Coulomb energies is varied in experiments either by changing both the magnetic field,  $B$ , and the electron density,  $n_s$ , at fixed filling factor,  $\nu = n_s \hbar c/eB$ , or by introducing a parallel component of the magnetic field. Ground-state spin transitions for the 2D electrons in GaAs have been observed using activation energy measurements at filling factor  $\nu = 2/3, 4/3, 3/5$ , etc. [1–3]; optical spectroscopy studies at  $\nu = 2/3, 2/5, 3/5$ , etc. [4]; nuclear magnetic resonance measurements at  $\nu = 2/3$  [5]; optical absorption experiments at  $\nu = 2/3$  [6]. The spin transitions occur at very different magnetic fields for different samples. For  $\nu = 2/3$ , the transition observed in Refs. [1–4] is the case at relatively low fields 2–4 T, while the transition found in Refs. [5, 6] occurs in considerably higher fields, about 8 T or yet higher. In contrast to the  $\nu = 2/3$  state, the quantum Hall state at  $\nu = 1/3$  is found to be fully spin polarized for all magnetic fields [6]. Still, spin-flip excitations at  $\nu = 1/3$  have been observed in transport [7], optical spectroscopy [8], and inelastic light scattering measurements [9].

Numerical calculations based on the composite fermion model are in qualitative agreement with the experimental results on the spin transitions [10, 11]. Within the model, the ground-state spin transition occurs when the difference in the Coulomb energies between the two ground states is balanced by the change in the Zeeman energy

due to spin polarization. The calculations for  $\nu = 2/3, 3/5$ , and  $4/7$  [11] underestimate the critical Zeeman energy that determines the magnetic field of the spin transition,  $B^*$ . In particular, these predict  $B^* \approx 1.5$  T for  $\nu = 2/3$ , assuming that the  $g$  factor is equal to  $|g| = 0.44$ . The discrepancy between theory and experiment cannot be explained by the finite thickness correction that is caused by finite extent of the wave functions in the direction perpendicular to the interface and gives rise to a decrease in the Coulomb energy [12, 13].

In this paper, we employ a magnetocapacitance technique to determine the magnetic field of the spin transition  $B^*$  at filling factor  $\nu = 2/3$ . The ratio of the Zeeman and Coulomb energies is varied in the experiment by a change of the back gate voltage that controls the wave function extent through the steepness of the confining potential in the direction perpendicular to the interface. As the wave function extent is increased, the field  $B^*$  is found to decrease appreciably. The calculated  $g$  factor change because of nonparabolicity effects and change of the Coulomb energy make approximately equal contributions to the shift of the spin transition field. Both mechanisms describe the observed relative shift of  $B^*$  with no fitting parameters. We check that the  $\nu = 2/3$  gap at fields above  $B^*$  increases with increasing wave function extent and reaches the value of the  $\nu = 1/3$  gap, indicating the gap symmetry in the high-field limit over the entire range of back gate voltages studied. The behavior of the  $\nu = 2/3$  gap clearly indicates the importance of spin effects.

Measurements were made in an Oxford dilution refrigerator with a base temperature of  $\approx 30$  mK on remotely doped GaAs/Al<sub>x</sub>Ga<sub>1-x</sub>As ( $x = 0.336$ ) single heterojunctions (with a low-temperature mobility  $\approx 4 \times 10^6$  cm<sup>2</sup>/Vs at electron density  $9 \times 10^{10}$  cm<sup>-2</sup>) having the quasi-Corbino geometry with area  $1.8 \times 10^4$   $\mu\text{m}^2$ . The depth of the 2D electron layer was 200 nm. A highly doped ( $1 \times 10^{18}$  cm<sup>-3</sup> Si) layer with thickness 100 nm was buried

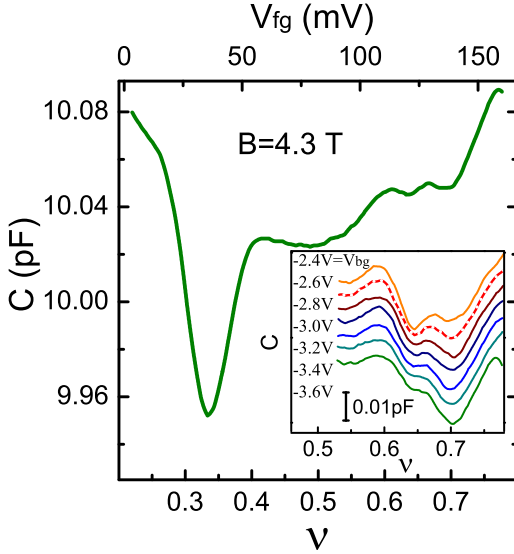


FIG. 1: Magnetocapacitance as a function of front gate voltage converted into filling factor in  $B = 4.3$  T at  $V_{bg} = -3.6$  V and  $T = 60$  mK. The inset shows an expanded view on  $C(\nu)$  near  $\nu = 2/3$  at different back gate voltages. The dashed line corresponds to the transition. Linear contribution to  $C(\nu)$  is subtracted and the traces are vertically shifted for clarity.

in the bulk of GaAs, at a distance  $5 \mu\text{m}$  from the interface. This layer remained well-conducting at low temperatures and served as a back electrode. A 200 nm low-temperature grown GaAs (LT-GaAs) layer between the back gate and the 2D electron layer was used to block leakage currents. The Fermi level in this layer is pinned near the midgap, which results in the formation of a Schottky barrier between  $n$ -GaAs and LT-GaAs [14]. A metallic front gate was deposited onto the surface of the sample. The presence of gates allowed variation of both the electron density and the confining potential by applying a dc bias between the gate and the 2D electrons. The front gate voltage was modulated with a small ac voltage of 2.5 mV at frequencies in the range 0.6–21 Hz, and both the imaginary and real components of the current were measured with high precision ( $\sim 10^{-16}$  A) using a current-voltage converter and a lock-in amplifier. Smallness of the real current component as well as proportionality of the imaginary current component to the excitation frequency ensure that we reach the low-frequency limit and the measured magnetocapacitance is not distorted by lateral transport effects. A dip in the magnetocapacitance in the quantum Hall effect is directly related to a jump of the chemical potential across a corresponding gap in the spectrum of the 2D electron system [15]:

$$\frac{1}{C} = \frac{1}{C_0} + \frac{1}{Ae^2 dn_s/d\mu}, \quad (1)$$

where  $C_0$  is the geometric capacitance between the gate and the 2D electrons,  $A$  is the sample area, and the derivative  $dn_s/d\mu$  of the electron density over the chemical potential is the thermodynamic density of states.

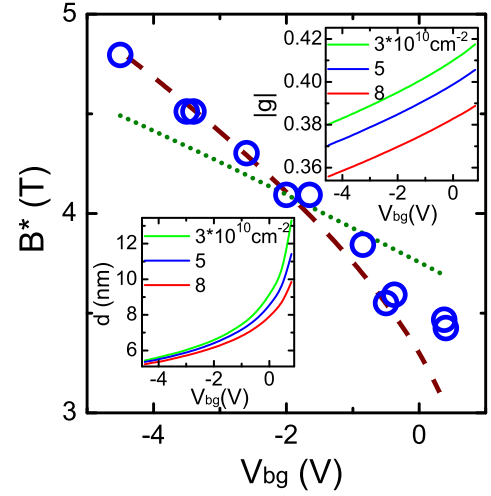


FIG. 2: Dependence of the magnetic field of the spin transition at  $\nu = 2/3$  on back gate voltage. The symbol size reflects the experimental uncertainty. The field  $B^*$  versus  $V_{bg}$  is calculated using Eq. (3) with  $F = 1$  (dotted green line) and the factor  $F$  taken from Ref. [11] (dashed brown line), where the coefficient  $\alpha$  is chosen so that  $B^*$  at  $V_{bg} = -2$  V corresponds to its experimental value. The calculated dispersion of the electron density distributions and single-electron  $g$  factor versus  $V_{bg}$  are displayed in the insets for different electron densities.

Near the filling factor  $\nu = 1/2$ , the capacitance  $C$  in the range of magnetic fields studied reaches its high-field value determined by the geometric capacitance  $C_0$ . The chemical potential jump  $\Delta\mu$  for electrons at fractional filling factor is determined by integrating the magnetocapacitance over the dip (for more details, see Ref. [16]):

$$\Delta\mu = \frac{e}{C_0} \int_{\text{dip}} (C_0 - C) dV_{fg}. \quad (2)$$

Note that the excitation gap corresponds to the chemical potential discontinuity divided by the fraction denominator [17].

A magnetocapacitance trace  $C$  as a function of front gate voltage  $V_{fg}$  converted into filling factor  $\nu$  is displayed in Fig. 1 for a magnetic field of 4.3 T and back gate voltage  $V_{bg} = -3.6$  V. A narrow dip and a double minimum in  $C$  are seen at  $\nu = 1/3$  and  $\nu = 2/3$ , respectively. The double minimum feature in  $C(\nu)$  at different back gate voltages is shown in the inset to Fig. 1. At fixed electron density (or magnetic field), the change of the confining potential is caused by a change of  $V_{bg}$  that is connected to a shift of  $V_{fg}$  via  $\Delta V_{fg} \approx -0.05 \Delta V_{bg}$ . One can see from the figure that the two minima reveal an interplay with changing back gate voltage. The interplay reflects the competition of the domains of two ground states at the critical point [18], similar to the case of the integer quantum Hall effect [19–21]. This corresponds to the spin transition in the ground state [16]. We determine the transition point as a point at which the two minima are approximately equal to each other (the dashed line

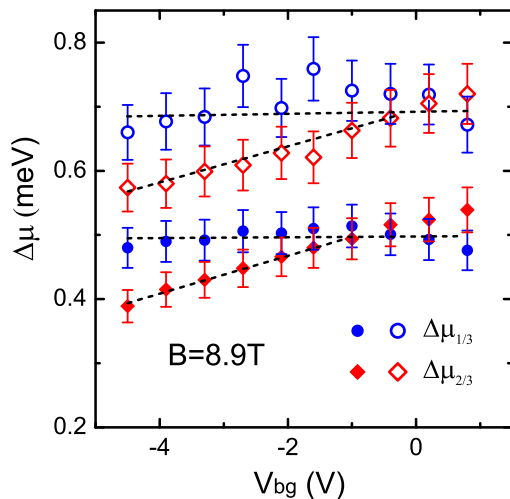


FIG. 3: The chemical potential jump at  $\nu = 1/3$  and  $\nu = 2/3$  versus back gate voltage for  $B = 8.9$  T and  $T = 0.2$  K (solid symbols). Also shown by the open symbols is the corrected data for  $\Delta\mu$ , see text. The dashed lines are guides to the eye.

in the figure).

In Fig. 2, we plot the field  $B^*$  of the spin transition for  $\nu = 2/3$  versus back gate voltage. The experimental value  $B^*$  decreases approximately linearly with increasing back gate voltage and remains well above the theoretical prediction  $B^* \approx 1.5$  T. Apparently, the increase of  $V_{bg}$  results in decreasing the steepness of the confining potential and increasing the wave function extent. Based on the competition between the Coulomb and Zeeman energies, one can tentatively expect that the reduction in the Coulomb energy due to the finite thickness of the 2D electron layer is responsible for the observed shift of the spin transition. However, the absolute value of  $g$  factor should increase with increasing wave function extent, which gives an alternative account of the shift of  $B^*$ .

We verify the importance of spin effects by measuring the chemical potential jump at  $\nu = 1/3$  and  $\nu = 2/3$  in magnetic fields above  $B^*$  as a function of back gate voltage, as shown by the solid symbols in Fig. 3. The data correspond to the limit of low temperatures where  $\Delta\mu$  saturates and becomes independent of temperature [16]. Whereas the gap at  $\nu = 1/3$  remains approximately constant, the  $\nu = 2/3$  gap increases with back gate voltage and reaches the value of the  $\nu = 1/3$  gap. The different behavior of the gaps signals the presence of spin-dependent contribution to the value of gap at  $\nu = 2/3$ , i.e., the increase of the  $\nu = 2/3$  gap with  $V_{bg}$  should be related to the increase in the absolute value of  $g$  factor. We do not observe in the experiment a decrease of the gaps with increasing back gate voltage, as expected from the suppression of gaps due to the finite thickness correction. This can be attributed to the manifestation of disorder effects: as the 2D electrons are pushed closer to scatterers at the interface, the gap suppression caused by disorder becomes stronger, compensating roughly for the finite

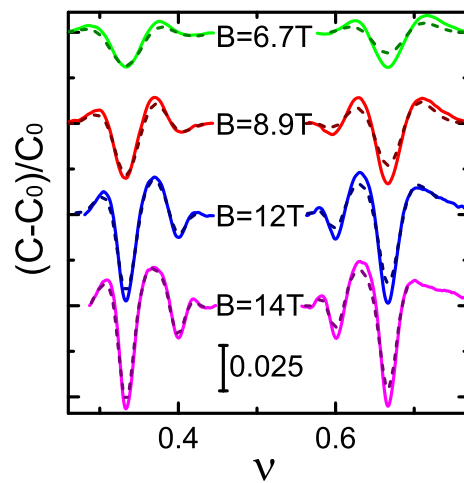


FIG. 4: Magnetocapacitance  $(C - C_0)/C_0$  as a function of filling factor for  $V_{bg} = -4.5$  V (dashed lines) and  $0.8$  V (solid lines) at different magnetic fields and temperatures 0.2, 0.2, 0.25, and 0.3 K (top to bottom). The curves are vertically shifted for clarity.

thickness correction. It is possible to take into account the effect of long-range disorder potential which leads to broadening the chemical potential jump as a function of filling factor. In the spirit of Ref. [22], linear extrapolations of the dependence  $\mu(\nu)$  to the fractional filling factor yield  $\Delta\mu$  for an ideal/homogeneous 2D electron system. The corrected data for  $\Delta\mu$  is shown by the open symbols in Fig. 3. One can see an upward shift in the dependences  $\Delta\mu(V_{bg})$  for both  $\nu = 1/3$  and  $\nu = 2/3$  so that their behavior remains basically the same. Therefore, it is the effect of short-range disorder potential that can be responsible for the compensation of the finite thickness correction.

Figure 4 displays the magnetocapacitance  $(C - C_0)/C_0$  versus  $\nu$  measured in the low-temperature limit at  $V_{bg} = -4.5$  V (dashed lines) and  $0.8$  V (solid lines) in different magnetic fields. As  $B$  is increased, the minima at  $\nu = 1/3$  and  $\nu = 2/3$  become symmetric and the value of the  $\nu = 2/3$  gap reaches that of the  $\nu = 1/3$  gap. Thus, the gap symmetry is the case at the highest magnetic fields over the entire range of back gate voltages studied. This implies that the spin-dependent contribution to the  $\nu = 2/3$  gap changes for the Coulomb contribution.

We would like to emphasize the significant difference between the theoretical descriptions of the ground-state spin transition and fractional gaps. The experimental values of fractional gaps are known to be noticeably smaller than the theoretical predictions, which can be attributed to disorder effects. In contrast, the magnetic fields of the spin transition measured in experiment exceed by far the theoretical values, which cannot be explained by the influence of disorder. Also, the  $g$  factor related to the spin transition is equal to its single-electron value, whereas the gap-related  $g$  factor can be enhanced due to interactions. Therefore, the case of spin transi-

tion includes less parameters and allows more rigorous comparison of experiment with theory.

It is easy to calculate the behavior of the single-electron  $g$  factor with back gate voltage for the samples used. By solving the Schrodinger and Poisson equations self-consistently [23], one obtains the dispersion,  $d$ , of the electron density distributions perpendicular to the interface as a function of back gate voltage for different  $n_s$  (bottom inset to Fig. 2). The value  $d$  gives a measure of the wave function extent and corresponds to the 2D subband bottom. This increases sharply with back gate voltage at high  $V_{bg}$ , close to the voltage  $V_{bg} \approx 0.9$  V for our samples, where the Schottky barrier between  $n$ -GaAs and LT-GaAs vanishes and the gate leakage current arises. In our case the variation of the  $g$  factor caused by wave function penetration in the AlGaAs barrier is small, and the dominant  $g$  factor change relative to  $g = -0.44$  in bulk GaAs originates from nonparabolicity effects. Within the  $kp$ -theory, the change at the 2D subband bottom is equal to  $\Delta g = \langle T \rangle / W$ , where  $\langle T \rangle$  is the average kinetic energy and  $W$  is the characteristic energy [24]. Using the experimental value of  $W \approx 150$  meV [25, 26] and taking account of the diamagnetic shift by half the cyclotron energy, we get the dependence of  $g$  on back gate voltage at different  $n_s$ , shown in the top inset of Fig. 2.

The behavior of the field  $B^*$  with back gate voltage is determined from the relation [11]

$$\alpha F \frac{e^2}{\kappa l_B} = \frac{1}{2} g \mu_B B, \quad (3)$$

where the coefficient  $\alpha$  is given by the difference in the Coulomb energies of fully and partially polarized states and the finite thickness correction  $F = 1$  ( $F < 1$ ) for zero (nonzero) thickness of the 2D electron layer. Assuming that  $F = 1$  and choosing the coefficient  $\alpha$  so that  $B^*$  at  $V_{bg} = -2$  V corresponds to its experimental value, we determine the relative shift of  $B^*$  caused by the  $g$  factor change (the dotted green line in Fig. 2). One can see in the figure that this mechanism accounts approximately for half of the observed shift of the magnetic field of the spin transition.

We now take into account the suppression of the Coulomb interactions due to the finite thickness of the 2D electron layer. Replacing the calculated wave function by the Fang-Howard wave function  $\zeta(z) = (b^3/2)^{1/2} z \exp(-bz/2)$  (where the  $z$ -axis is perpendicular to the interface,  $z > 0$  in substrate, and  $b$  is the parameter) [27] with the same dispersion  $d$ , one determines the value  $1/bl_B = d/\sqrt{3}l_B$  that controls the suppression factor  $F$ . Using the dependence  $F(1/bl_B)$  calculated in

Ref. [11], we obtain from Eq. (3) the total relative shift of  $B^*$  (the dashed brown line in Fig. 2), which is in agreement with the experiment [28]. Thus, both mechanisms make approximately equal contributions to the change of the spin transition field and describe the observed relative shift of  $B^*$  with no fitting parameters.

The theoretical magnetic field of the spin transition [11] is approximately two times smaller than the value observed in the experiment. That is to say, the theory in question underestimates by  $\approx 30\%$  the difference in the Coulomb energies between fully and partially polarized states, which is experimentally equal to  $\approx 0.005e^2/\kappa l_B \approx 0.04$  meV. The discrepancy with the experiment may result from distant extrapolations used to determine the Coulomb energies. Note that the aforementioned significant difference in the fields  $B^*$  for different samples [1–6] as well as the lack of symmetry between the  $\nu = 1/3$  and  $\nu = 2/3$  gaps in the same magnetic field (see, e.g., Ref. [18]) are likely to be caused mainly by the reduced absolute values of  $g$  factor for the particular sample design.

The gap-related  $g$  factor can be determined assuming that the sum of the gaps at  $\nu = 1/3$  and  $\nu = 2/3$  is equal to the Zeeman energy. Using the corrected data in Fig. 3, one estimates an enhanced  $g$  factor at  $g \approx 0.9$ . Its variation  $\delta g \approx 0.1$  in the range of back gate voltages studied is bigger than the change of the single-electron  $g$  factor (top inset of Fig. 2). This can be explained qualitatively by the many-body enhancement of the  $g$  factor, depending on the disorder and finite thickness effects.

In summary, we have found that the magnetic field of the spin transition at filling factor  $\nu = 2/3$  decreases appreciably as the wave function extent controlled by back gate voltage is increased. Our calculations show that the contributions to the shift of  $B^*$  from the  $g$  factor change because of nonparabolicity and the change of the Coulomb energy are approximately the same. The observed relative shift of  $B^*$  is described with no fitting parameters. We have checked that the  $\nu = 2/3$  gap at fields above  $B^*$  increases with increasing wave function extent and reaches the value of the  $\nu = 1/3$  gap, indicating the gap symmetry in the high-field limit over the entire range of back gate voltages studied. The behavior of the  $\nu = 2/3$  gap signals the importance of spin effects.

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